Generation of energy selective excitations in quantum Hall edge states

C. Leicht,[1] P. Mirovsky,[1] B. Kaestner,[1] F. Hohls,[1] V. Kashcheyevs,[2,3] E. V. Kurganova,[4] U. Zeitler,[4] T. Weimann,[1] K. Pierz,[1] and H. W. Schumacher[1]

[1]*Physikalisch-Technische Bundesanstalt, 38116 Braunschweig, Germany.*

[2]*Faculty of Physics and Mathematics, University of Latvia, Riga LV-1002 Latvia.*

[3]*Faculty of Computing, University of Latvia, Riga LV-1586, Latvia.*

[4]*High Field Magnet Laboratory, Institute for Molecules and Materials, Radboud University Nijmegen, 6525 ED Nijmegen, The Netherlands.*

We operate an on-demand source of single electrons in high perpendicular magnetic fields up to 30 T, corresponding to a filling factor ν below 1/3. The device extracts and emits single charges at a tunable energy from and to a two-dimensional electron gas, brought into well defined integer and fractional quantum Hall (QH) states. It can therefore be used for sensitive electrical transport studies, e.g. of excitations and relaxation processes in QH edge states.

Charge transport in two-dimensional electron gases placed in a strong perpendicular magnetic field is ruled by chiral edge states.[1–3] These edge states are now beeing exploited routinely in fundamental physics experiments, e.g. in electron interferometers.[4,5] Moreover, gapless neutral edge excitations have been predicted,[6,7] though not yet directly observed in experiments using quantum point contacts to generate edge excitations, as performed in e.g. Refs. 8 and 9. Additional counterpropagating edge excitations in the fractional quantum Hall (QH) state have also been predicted,[3,10] but were not found in studies of edge magneto plasmons.[11,12] Only very recently a shot noise experiment found first indications for a neutral counterpropagating mode.[13]

In this paper we demonstrate a new method to generate triggered single energy selective excitations in integer and fractional QH edges to probe possible edge excitations and relaxation processes. Furthermore, this method allows to precisely control the emission statistics of the electrons, which opens the possibility for efficient time resolved measurements.

We adapt a structure that has previously been employed as high precision current source, both in the dc[14] and ac regime.[15] A schematic of our device and an electron micrograph are shown in Fig. 1a. It was realized in an AlGaAs/GaAs heterostructure. A 700 nm wide constriction was wet-etched inside a two-dimensional electron gas. The device was contacted at source (S) and drain (D) using an annealed layer of AuGeNi. The constriction is crossed by Ti-Au finger gates G_1 and G_2. A quantum dot (QD) with a quasibound state ψ is formed by applying voltages V_1 and V_2 to G_1 and G_2, respectively; a third gate G_3 is not used and set to ground. The corresponding potential landscape along the constriction is shown in Fig. 1b. An additional sinusoidal signal of power P^{RF} and frequency f is coupled to G_1 and varies both the height of the barrier and the energy $\varepsilon(t) = \varepsilon_1 \cos\omega t + \varepsilon_0$ of the quasibound state ($\omega = 2\pi f$). During the first half cycle $\varepsilon(t)$ drops below the chemical potential μ_S and ψ is loaded with an electron with energy $\mu_S - E_L$ [see Fig. 1(b)]. During the second half-cycle, $\varepsilon(t)$ is raised sufficiently fast above μ_D and the electron can

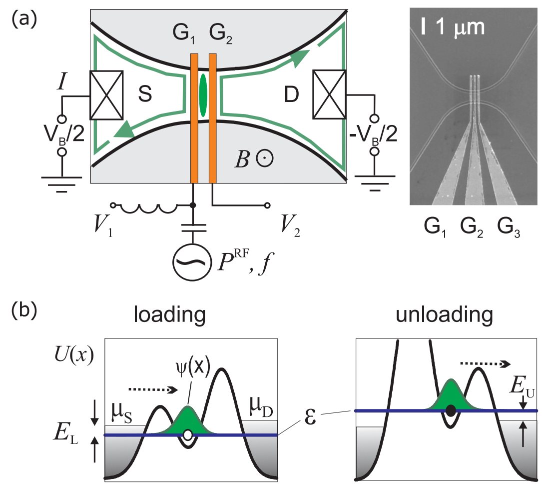

FIG. 1. (color online) Description of the device and operating principle. (a) Left, schematic of the device. Electron micrograph of the sample shown on the right. (b) Schematic of the potential energy landscape along the channel for the stages of loading and unloading.

be unloaded to the drain with an excess energy E_U. This process, resulting into a *quantized* current $I = e \cdot f$ with e the electron charge, is non-adiabatic and requires that the loaded QD state is raised sufficiently fast through the chemical potentials $\mu_{S/D}$ to avoid unwanted charge transfer.[14] The scheme can be generalized to a quantized transport of n electrons per cycle, i.e. $I = n \cdot e \cdot f$, where n can be derived from the decay cascade model.[16]

The current is accompanied by a periodic excitation in the drain at energy E_U above μ_D. Upon application of a strong perpendicular magnetic field B, transport in S and D takes place via edge channels, marked symbolically with green arrows in Fig. 1a. Using the *dynamical* QD it is now possible to trigger single energy selective excitation quanta of the QH edge state.

The number of electrons emitted into D per cycle may be tuned using V_2, as shown in Fig. 2 for a measurement carried out in a ^{3}He cryostat with a base temperature of 300 mK. Under zero-field conditions approximately one

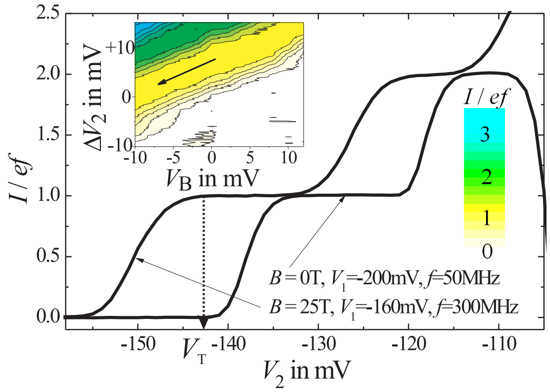

FIG. 2. (color online) Normalized current I/ef generated by electrons emitted into drain as function of V_2. The quantized regime $I = ef$ under $B = 0$ and $B = 25\,\mathrm{T}$ conditions is obtained over a V_2 range of several mV. The threshold voltage V_T is indicated by the dashed arrow, beyond which the quantized regime breaks down. Inset showing I/ef as the bias V_B is varied.

electron is emitted per cycle to D for $V_2 = -135 \cdots -120\,\mathrm{mV}$. If a perpendicular magnetic field is applied it has been found previously in acoustically driven dynamical QDs that quantization is quenched for $B \geq 1\,\mathrm{T}$.[17] In the present case, where the dynamic potential is generated directly by gates, quantization can be achieved up to very high magnetic fields, as shown in Fig. 2. Here emission of single charges into D, i.e. quantized charge pumping, at $B = 25\,\mathrm{T}$ is shown, corresponding to a fractional filling factor $\nu = 1/3$ of the undisturbed QH liquid.

The emission energy E_U depends on f as well as on the bias voltage, V_B, and V_2. It may be determined experimentally using gates as energy filter, as for instance used in Ref. 18. We have obtained a first estimate for $\Delta E = E_U + E_L$ at $B = 0\,\mathrm{T}$ based on the variation of the ef-plateau lengths along V_1 as function of P^{RF}.[19] For the studied device $\Delta E \approx 14$ to $17\,\mathrm{meV}$ for frequencies ranging from $50 \cdots 300\,\mathrm{MHz}$, and for V_2 set to the negative side of the plateau ($V_2 = V_T$, see Fig. 2).

In the following we estimate the distribution of the energy of the emitted electrons, $p\,(E_U)$, based on the Master equation model of Ref. 14. The sharpness of the distribution, ΔE_U, can be tuned by the *selectivity* $s \equiv g/\varepsilon_1$ of the barrier at G_2 with $g \equiv \ln \Gamma_2^{\max}/\Gamma_2^{\min}$. Here $\Gamma_2^{\max/\min}$ are the maximal and minimal tunnel rates during one cycle of modulation, where we also assume that Γ_2 depends exponentially on ε. To obtain an expression for the energy distribution we consider the case when unloading ($\varepsilon \geq \mu_D$) takes place during the phase when ε changes most rapidly, such that $\varepsilon(t) \approx \varepsilon_1 \omega t + \mu_D$. The problem can then be simplifed to a dynamic QD completely occupied at $t \to -\infty$ which unloads to drain via G_2 with increasing rate $\Gamma_2(t) = \Gamma_2^0 \exp\left(\frac{1}{2}\omega\, g\, t\right)$, where Γ_2^0 is the escape rate when $\varepsilon(t) = \mu_D$. With these assumptions the distribution of the emission times is peaked at $t_e = \beta^{-1} \ln\left(\beta/\Gamma_2^0\right)$ with $\beta \equiv g\,\omega/2$. The width of the corresponding energy distribution is then given by $\Delta E_U = 2\,/s$. Hence, to obtain a narrow emission energy distribution one may optimize the barrier shape of G_2 to maximize s. The lowest achievable ΔE_U is limited by the quantum-mechanical uncertainty of energy, on the order of $\hbar\Gamma_2(t_e) = (g/2)\hbar\omega$. For the frequencies chosen in this experiment the minimal ΔE_U lies in the μeV range.

The derivation above also shows that the emission energy $E_U = \varepsilon(t_e) - \mu_D$ depends on the frequency ω and the tunnel rate Γ_2^0 logarithmically,

$$E_U \approx \varepsilon_1 \omega t_e = \Delta E_U \, \ln\left(\frac{\varepsilon_1\,\omega}{\Delta E_U\,\Gamma_2^0}\right). \qquad (1)$$

Since typically Γ_2^0 depends on V_2 exponentially, the gate voltage can be readily used to tune the emission energy, i.e. $E_U \propto -|e|\,V_2$. To ensure single triggered excitation events (within a certain error margin) V_2 may be tuned only within the plateau voltage range, i.e. where $I \approx e\,f$. The highest energy is obtained for the transition voltage, V_T, where $I = ef$ switches to $I = 0$, i.e. close to the negative side of the plateau where Γ_2^0 is minimal (see Fig. 2). From Eq. 1 it follows that increasing the modulation amplitude ε_1 enhances E_U only logarithmically. To extend the energy range efficiently, the bias voltage $V_B \equiv (\mu_D - \mu_S)\,/|e|$ may be made more negative, decreasing Γ_2^0 since the condition $\varepsilon = \mu_D$ will then take place earlier in the cycle, i.e. $E_U \propto -|e|\,V_B$. At the same time the chance for emitting an additional electron increases, as seen from the inset in Fig. 2. This behaviour is consistent with the decay cascade model,[16] considering that the time t_c at which the decay cascade starts is given by $\varepsilon(t_c) \equiv \mu_S$. The corresponding escape rate at G_1, $\Gamma_1(\varepsilon(t_c))$, controls the number of electrons captured per cycle. To remain in the quantized regime, V_2 and consequently Γ_2^0 have to be decreased as indicated by the arrow in the inset of Fig. 2, leading to an additional enhancement of E_U according to Eq. 1. Hence, combining the f -, V_B - and V_2 - dependence an excitation energy range up to several tens of meV should be possible using this technique. Despite the potentially large energy, the heating of the edge state can be kept at a minimum by choosing a sufficiently low frequency. Furthermore, this energy selective and time controlled excitation source could be combined with selective edge mode detection[20] and a time-gated detector technique[12] for sensitive studies of the underlying transport processes.

For the presented excitation source we require that the gates G_1 and G_2 coincide with the border of the undisturbed QH liquid, in order to avoid broadening of the energy distribution $p(E_U)$. In previous studies of this dynamical QD in perpendicular magentic field, such as in Refs. 21 and 22, the electron density of states and the corresponding filling factor of the leads connecting to the QD via G_1 and G_2 could not be established. In those works a wire of constant nominal width was employed, where side wall depletion may result in varying electron densities inside the wire, different from the undisturbed

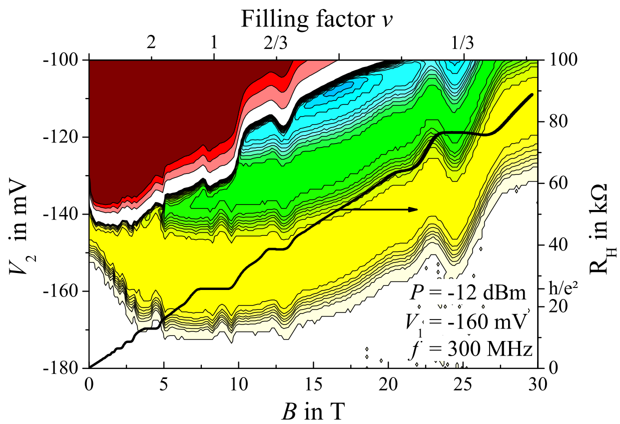

FIG. 3. (color online) Normalized current I/ef as a function of V_2 and B, as well as Hall resistance R_H as a function of B. The power and frequency have been chosen to remain in the decay cascade regime.[16] The colors and contours in the diagramme correspond to those in Fig. 2.

QH liquid. The tapered channel geometry used in the present work intends to avoid this complication. Fig. 3 shows evidence that in the case of the tapered channel shape the QD extracts and emits electrons directly from and to the undisturbed QH liquid. We conclude this from the oscillations in $V_T(B)$ which coincide with the superimposed Hall resistance R_H determined for the undisturbed QH liquid. We relate these oscillations to the variations in $\mu_{S/D}$ for transitions between different integer and fractional filling factors[23] modifying the decay rates which V_T is sensitive to. In particular, the data in Fig. 3 demonstrate the clocked capturing of electrons directly from a fractional QH edge state.

Similar oscillations at lower magnetic fields have also been reported in Refs. 17 and 22. In Ref. 17 quantization was quenched for $B \geq 1T$ and no clear comparison seems possible. The periodicity observed in Ref. 22 does not correspond to the R_H variation inferred from the charge carrier density specified. This observation indicates emission into a localized region of reduced electron density inside the etched channel.

Finally we note that charge pumping from fractional edge states as demonstrated here may be developed further into the realization of a fractional charge pump as proposed by Simon,[24] which may be used as a measurement of the charge of the fractional quantum Hall quasiparticle.

We thank S. J. Wright, M. Kataoka, S. P. Giblin and J. D. Fletcher for useful discussions. Technical support from H. Marx and U. Becker is gratefully acknowledged. This work has been supported by ERANET Plus under EU Grant Agreement No. 217257. C.L. has been supported by International Graduate School of Metrology, Braunschweig. Part of this work has been supported by EuroMagNET II under the EU contract number 228043.

REFERENCES

[1] B. I. Halperin, Phys. Rev. B, **25**, 2185 (1982).
[2] M. Büttiker, Phys. Rev. B, **38**, 9375 (1988).
[3] X. G. Wen, Phys. Rev. Lett., **64**, 2206 (1990).
[4] Y. Ji, Y. Chung, D. Sprinzak, M. Heiblum, D. Mahalu, and H. Shtrikman, Nature, **422**, 415 (2003).
[5] I. Neder, N. Ofek, Y. Chung, M. Heiblum, D. Mahalu, and V. Umansky, Nature, **448**, 333 (2007).
[6] I. L. Aleiner and L. I. Glazman, Phys. Rev. Lett., **72**, 2935 (1994).
[7] C. d. C. Chamon and X. G. Wen, Phys. Rev. B, **49**, 8227 (1994).
[8] C. Altimiras, H. le Sueur, U. Gennser, A. Cavanna, D. Mailly, and F. Pierre, Nature Physics, **6**, 34 (2010).
[9] G. Granger, J. P. Eisenstein, and J. L. Reno, Phys. Rev. Lett., **102**, 086803 (2009).
[10] A. H. MacDonald, Phys. Rev. Lett., **64**, 220 (1990).
[11] R. C. Ashoori, H. L. Stormer, L. N. Pfeiffer, K. W. Baldwin, and K. West, Phys. Rev. B, **45**, 3894 (1992).
[12] G. Ernst, N. B. Zhitenev, R. J. Haug, and K. von Klitzing, Phys. Rev. Lett., **79**, 3748 (1997).
[13] A. Bid, N. Ofek, M. Heiblum, V. Umansky, and D. Mahalu, Phys. Rev. Lett., **103**, 236802 (2009).
[14] B. Kaestner, V. Kashcheyevs, S. Amakawa, M. D. Blumenthal, L. Li, T. J. B. M. Janssen, G. Hein, K. Pierz, T. Weimann, U. Siegner, and H. W. Schumacher, Phys. Rev. B, **77**, 153301 (2008),.
[15] G. Feve, A. Mahe, J.-M. Berroir, T. Kontos, B. Placais, D. C. Glattli, A. Cavvanna, B. Etienne, and Y. Jin, Science, **316**, 1169 (2007).
[16] V. Kashcheyevs and B. Kaestner, Phys. Rev. Lett., **104**, 186805 (2010).
[17] J. Cunningham, V. I. Talyanskii, J. M. Shilton, M. Pepper, A. Kristensen, and P. E. Lindelof, Phys. Rev. B, **62**, 1564 (2000).
[18] F. Hohls, M. Pepper, J. P. Griffiths, G. A. C. Jones, and D. A. Ritchie, Appl. Phys. Lett., **89**, 212103 (2006).
[19] B. Kaestner, V. Kashcheyevs, G. Hein, K. Pierz, U. Siegner, and H. W. Schumacher, Appl. Phys. Lett., **92**, 192106 (2008).
[20] G. Sukhodub, F. Hohls, and R. J. Haug, Phys. Rev. Lett., **93**, 196801 (2004).
[21] B. Kaestner, C. Leicht, V. Kashcheyevs, K. Pierz, U. Siegner, and H. W. Schumacher, Appl. Phys. Lett., **94**, 012106 (2009).
[22] S. J. Wright, M. D. Blumenthal, G. Gumbs, A. L. Thorn, M. Pepper, T. J. B. M. Janssen, S. N. Holmes, D. Anderson, G. A. C. Jones, C. A. Nicoll, and D. A. Ritchie, Phys. Rev. B, **78**, 233311 (2008).
[23] R. E. Prange and S. M. Girvin, eds., *The Quantum Hall Effect* (Springer Verlag, 1986).
[24] S. H. Simon, Phys. Rev. B, **61**, R16327 (2000).